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Diffractive light quark jet production at hadron colliders in the two-gluon exchange model

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Abstract

Massless quark and antiquark jet production at large transverse momentum in the coherent diffractive processes at hadron colliders is calculated in the two-gluon exchange parametrization of the Pomeron model. We use the helicity amplitude method to calculate the cross section formula. We find that for the light quark jet production the diffractive process is related to the differential off-diagonal gluon distribution function in the proton. We estimate the production rate for this process at the Fermilab Tevatron by approximating the off-diagonal gluon distribution function by the usual diagonal gluon distribution in the proton. And we find that the cross sections for the diffractive light quark jet production and the charm quark jet production are in the same order of magnitude. We also use the helicity amplitude method to calculate the diffractive charm jet production at hadron colliders, by which we reproduce the leading logarithmic approximation result of this process we previously calculated.

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I. INTRODUCTION

In recent years, there has been a renaissance of interest in diffractive scattering. These diffractive processes are described by the Regge theory in terms of the Pomeron ($I\!P$) exchange [1]. The Pomeron carries quantum numbers of the vacuum, so it is a colorless entity in QCD language, which may lead to the "rapidity gap" events in experiments. However, the nature of Pomeron and its reaction with hadrons remain a mystery. For a long time it had been understood that the dynamics of the "soft pomeron" is deeply tied to confinement. However, it has been realized now that how much can be learned about QCD from the wide variety of small-x and hard diffractive processes, which are now under study experimentally. In Refs. [2,3], the diffractive J/ψ and Υ production cross section have been formulated in photoproduction processes and in DIS processes in perturbative QCD. In the framework of perturbative QCD the Pomeron is represented by a pair of gluon in the color-singlet sate. This two-gluon exchange model can successfully describe the experimental results from HERA [5].

On the other hand, as we know that there exist nonfactorization effects in the hard diffractive processes at hadron colliders [6–9]. First, there is the so-called spectator effect [8], which can change the probability of the diffractive hadron emerging from collisions intact. Practically, a suppression factor (or survive factor) "S_F" is used to describe this effect [10]. Obviously, this suppression factor can not be calculated in perturbative QCD, which is now viewed as a nonperturbative parameter. Typically, the suppression factor S_F is determined to be about 0.1 at the energy scale of the Fermilab Tevatron [9]. Another nonfactorization effect discussed in literature is associated with the coherent diffractive processes at hadron colliders [7], in which the whole Pomeron is induced in the hard scattering. It is proved in [7] that the existence of the leading twist coherent diffractive processes is associated with a breakdown of the QCD factorization theorem.

Based on the success of the two-gluon exchange parametrization of the Pomeron model in the description of the diffractive photoproduction processes at ep colliders [2,3,5], we may extend the applications of this model to calculate the diffractive processes at hadron colliders in perturbative QCD. Under this context, the Pomeron represented by a color-singlet two-gluon system emits from one hadron and interacts with another hadron in hard process, in which the two gluons are both involved (as shown in Fig. 1). The partonic process is plotted in Fig.2, where there are nine diagrams in the leading order of perturbative QCD. Therefore, these processes calculated in the two-gluon exchange model are just belong to the coherent diffractive processes in hadron collisions. Another important feature of the calculations of the diffractive processes in this model recently demonstrated is the sensitivity to the off-diagonal parton distribution function in the proton [11].

Using this two-gluon exchange model, we have calculated the diffractive J/ψ production [12], charm jet production [13], massive muon pair and W boson productions [14] in hadron collisions. These calculations show that we can explore much low x (off-diagonal) gluon distribution function and study the coherent diffractive processes at hadron colliders through these processes. In this paper, we will calculate the light quark (massless) jet production at large transverse momentum in the coherent diffractive processes at hadron colliders by using the two-gluon exchange model. In the calculations of Refs. [12–14], there always is a large mass scale associated with the production process. That is M_ψ for J/ψ production,

m_c for the charm jet production, M^2 for the massive muon production (M^2 is the invariant mass of the muon pair) and M_W^2 for W boson production. However, in the light quark jet production process, there is no large mass scale. So, for the light quark jet production, the large transverse momentum is needed to guarantee the application of the perturbative QCD. Furthermore, the experience on the calculations of the diffractive di-quark photoproduction [15] shows that the light quark jet production in the two-gluon exchange model has a distinctive feature that there is no contribution from the small l_T^2 region ($l_T^2 < k_T^2$) in the integration of the amplitude over l_T^2. So, the expansion (in terms of l_T^2/M_X^2) method used in Refs. [12–14] can not be applied to the calculations of light quark jet production. In the following calculations, we will employ the helicity amplitude method to calculate the amplitude of the diffractive light quark jet production in hadron collisions. We will show that the production cross section is related to the differential (off-diagonal) gluon distribution function in the proton as that in the diffractive di-quark jet photoprodution process [15]. (On the other hand, we note that the cross sections of the processes calculated in Refs. [12–14] are related to the integrated gluon distribution function in the proton).

The diffractive production of heavy quark jet at hadron colliders has also been studied by using the two-gluon exchange model in Ref. [16]. However, their calculation method is very different from ours [1]. In their calculations, they separated their diagrams into two parts, and called one part the coherent diffractive contribution to the heavy quark production. However, this separation can not guarantee the gauge invariance [13]. In our approach, we follow the definition of Ref. [7], i.e., we call the process in which the whole Pomeron participants in the hard scattering process as the coherent diffractive process. Under this definition, all of the diagrams plotted in Fig.2 for the partonic process $gp \to q\bar{q}p$ contribute to the coherent diffractive production.

The rest of the paper is organized as follows. In Sec.II, we will give the cross section formula for the partonic process $gp \to q\bar{q}p$ in the leading order of perturbative QCD, where we employ the helicity amplitude method to calculate the amplitude for this process. In Sec.III, we use this helicity amplitude method to recalculate the diffractive charm jet production process $gp \to c\bar{c}p$, by which we will reproduce the leading logarithmic approximation result for this process previously calculated in Ref. [13]. In Sec.IV, we will estimate the production rate of diffractive light quark jet at the Fermilab Tevatron by approximating the off-diagonal gluon distribution function by the usual diagonal gluon distribution function in the proton. And the conclusions will be given in Sec.V.

II. THE CROSS SECTION FORMULA FOR THE PARTONIC PROCESS

For the partonic process $gp \to q\bar{q}p$, in the leading order of perturbative QCD, there are nine diagrams shown in Fig.2. The two-gluon system coupled to the proton (antiproton) in Fig.2 is in a color-singlet state, which characterizes the diffractive processes in perturbative QCD. Due to the positive signature of these diagrams (color-singlet exchange), we know that the real part of the amplitude cancels out in the leading logarithmic approximation.

[1]For detailed discussions and comments, please see [13]

To get the imaginary part of the amplitude, we must calculate the discontinuity represented by the crosses in each diagram of Fig.2.

The first four diagrams of Fig.2 are the same as those calculated in the diffractive photoproduction processes. But, due to the existence of gluon-gluon interaction vertex in QCD, in the partonic process $gp \to q\bar{q}p$, there are additional five diagrams (Fig.2(5)-(9)). These five diagrams are needed for complete calculations in this order of QCD.

In our calculations, we express the formulas in terms of the Sudakov variables. That is, every four-momenta k_i are decomposed as,

$$k_i = \alpha_i q + \beta_i p + \vec{k}_{iT}, \tag{1}$$

where q and p are the momenta of the incident gluon and the proton, $q^2 = 0$, $p^2 = 0$, and $2p \cdot q = W^2 = s$. Here s is the c.m. energy of the gluon-proton system, i.e., the invariant mass of the partonic process $gp \to q\bar{q}p$. α_i and β_i are the momentum fractions of q and p respectively. k_{iT} is the transverse momentum, which satisfies

$$k_{iT} \cdot q = 0, \quad k_{iT} \cdot p = 0. \tag{2}$$

All of the Sudakov variables for every momentum are determined by using the on-shell conditions of the momenta represented by the external lines and the crossed lines in the diagram. The calculations of these Sudakov variables are similar to those in the diffractive charm jet production process $gp \to c\bar{c}p$. And all of these Sudakov variables for the process $gp \to q\bar{q}p$ of the light quark jet production can take their values from the corresponding formulas in Ref. [13] after taking $m_c \to 0$. For convenience, we list all of the Sudakov variables for the diffractive process $gp \to q\bar{q}p$ in the following.

For the momentum u, we have

$$\alpha_u = 0, \quad \beta_u = x_{IP} = \frac{M_X^2}{s}, \quad u_T^2 = t = 0, \tag{3}$$

where M_X^2 is the invariant mass squared of the diffractive final state including the light quark and antiquark jets. For the high energy diffractive process, we know that $M_x^2 \ll s$, so we have β_u (x_{IP}) as a small parameter. For the momentum k,

$$\alpha_k(1+\alpha_k) = -\frac{k_T^2}{M_X^2}, \quad \beta_k = -\alpha_k \beta_u, \tag{4}$$

where k_T is the transverse momentum of the out going quark jet. For the loop mentum l,

$$\begin{aligned}
\alpha_l &= -\frac{l_T^2}{s}, \\
\beta_l &= \frac{2(k_T, l_T) - l_T^2}{\alpha_k s}, \quad \text{for Diag.1, 3, 5,} \\
&= \frac{2(k_T, l_T) + l_T^2}{(1+\alpha_k) s}, \quad \text{for Diag.2, 4, 6,} \\
&= -\frac{M_X^2 - l_T^2}{s}, \quad \text{for Diag.7, 8, 9.}
\end{aligned} \tag{5}$$

Using these Sudakov variables, we can give the cross section formula for the partonic process $gp \to q\bar{q}p$ as,

$$\frac{d\hat{\sigma}(gp \to q\bar{q}p)}{dt}|_{t=0} = \frac{dM_X^2 d^2k_T d\alpha_k}{16\pi s^2 16\pi^3 M_X^2}\delta(\alpha_k(1+\alpha_k)+\frac{k_T^2}{M_X^2})\sum\overline{|\mathcal{A}|}^2, \tag{6}$$

where $\mathcal{A}$ is the amplitude of the process $gp \to q\bar{q}p$. We know that the real part of the amplitude $\mathcal{A}$ is zero, and the imaginary part of the amplitude $\mathcal{A}(gp \to q\bar{q}p)$ for each diagram of Fig.2 has the following general form,

$$\mathrm{Im}\mathcal{A} = C_F(T_{ij}^a)\int\frac{d^2l_T}{(l_T^2)^2}F \times \bar{u}_i(k+q)\Gamma_\mu v_j(u-k), \tag{7}$$

where C_F is the color factor for each diagram, and is the same as that for the $gp \to c\bar{c}p$ process [13]. a is the color index of the incident gluon. Γ_μ represents some γ matrices including one propagator. F in the integral is defined as

$$F = \frac{3}{2s}g_s^3 f(x', x''; l_T^2), \tag{8}$$

where

$$f(x', x''; l_T^2) = \frac{\partial G(x', x''; l_T^2)}{\partial \mathrm{ln} l_T^2}, \tag{9}$$

where the function $G(x', x''; k_T^2)$ is the so-called off-diagonal gluon distribution function [11]. Here, x' and x'' are the momentum fractions of the proton carried by the two gluons. It is expected that at small x, there is no big difference between the off-diagonal and the usual diagonal gluon densities [17]. So, in the following calculations, we estimate the production rate by approximating the off-diagonal gluon density by the usual diagonal gluon density, $G(x', x''; Q^2) \approx xg(x, Q^2)$, where $x = x_{I\!P} = M_X^2/s$.

In Ref. [13], we calculate the amplitude (7) for the heavy quark diffractive production processes by expanding Γ_μ in terms of l_T^2. However, in the light quark jet production process $gp \to q\bar{q}p$, the expansion method is not yet valid. According to the result of [13], the production cross section is proportional to the heavy quark mass. If we apply this formula to the light quark jet production, the cross section will be zero. That is to say, the expansion of the amplitude in terms of l_T^2, in which the large logarithmic contribution comes from the region of $l_T^2 \ll M_X^2$, is not further suitable for the calculation of the cross section for massless light quark jet production. Furthermore, the experience of the calculation of the diffractive light quark jet photoproduction process $\gamma p \to q\bar{q}p$ (for $Q^2 = 0$) [15] indicates that there is no contribution from the region of $l_T^2 < k_T^2$ in the integration of the amplitude over l_T^2. In the hadroproduction process $gp \to q\bar{q}p$, the situation is the same. So, the expansion method, in which l_T^2 is taken as a small parameter, is not valid for the calculations of the massless quark production processes.

In the following, we employ the helicity amplitude method [18] to calculate the amplitude Eq.(7). For the massless quark spinors, we define

$$u_\pm(p) = \frac{1}{\sqrt{2}}(1 \pm \gamma_5)u(p). \tag{10}$$

For the polarization vector of the incident gluon, which is transversely polarized, we choose,

$$e_{\pm} = \frac{1}{\sqrt{2}}(0, 1, \pm i, 0). \tag{11}$$

The helicity amplitudes for the processes in which the polarized Dirac particles are involved have the following general forms [18],

$$\bar{u}_{\pm}(p_f)Qv_{\mp}(p_i) = \frac{Tr[Q \not{p}_i \not{n} \not{p}_f(1 \mp \gamma_5)]}{4\sqrt{(n \cdot p_i)(n \cdot p_f)}}, \tag{12}$$

where n is an arbitrary massless 4-vector, which is set to be $n = p$ in the following calculations. Using this formula (12), the calculations of the helicity amplitude $\mathcal{A}(\lambda(g), \lambda(q), \lambda(\bar{q}))$ for the diffractive process $gp \to q\bar{q}p$ is straightforward. Here λ are the corresponding helicities of the external gluon, quark and antiquark. In our calculations, we only take the leading order contribution, and neglect the higher order contributions which are proportional to $\beta_u = \frac{M_X^2}{s}$ because in the high energy diffractive processes we have $\beta_u \ll 1$.

For the first four diagrams, to sum up together, the imaginary part of the amplitude $\mathcal{A}(\pm, +, -)$ is

$$\mathrm{Im}\mathcal{A}^{1234}(\pm, +, -) = \alpha_k^2(1 + \alpha_k)\mathcal{N} \times \int \frac{d^2\vec{l}_T}{(l_T^2)^2} f(x', x''; l_T^2)(-\frac{2}{9}\frac{\vec{e}^{(\pm)} \cdot \vec{k}_T}{k_T^2} - \frac{1}{36}\frac{\vec{e}^{(\pm)} \cdot (\vec{k}_T - \vec{l}_T)}{(\vec{k}_T - \vec{l}_T)^2}), \tag{13}$$

where $\frac{2}{9}$ and $-\frac{1}{36}$ are the color factors for Diag.1,4 and Diag.2,3 respectively, and $\mathcal{N}$ is defined as

$$\mathcal{N} = \frac{s}{\sqrt{-\alpha_k(1 + \alpha_k)}} g_s^3 T_{ij}^a. \tag{14}$$

The other helicity amplitudes for the first four diagrams have the similar form as (13). The amplitude expression Eq. (13) is the same as that for the photoproduction process $\gamma p \to q\bar{q}p$ previously calculated in Refs. [4,15] except the difference on the color factors. In the diffractive photoproduction process, the color factors of these four diagrams are the same (they are all $\frac{2}{9}$), while in hadroproduction process the color factors are no longer the same for these four diagrams. It is instructive to see what is the consequence of this difference. We know that the amplitude of the diffractive process in Eq. (7) must be zero in the limit $l_T^2 \to 0$. Otherwise, this will lead to a linear singularity when we perform the integration of the amplitude over l_T^2 due to existence of the factor $1/(l_T^2)^2$ in the integral of Eq. (7) [13]. This linear singularity is not proper in QCD calculations. So, we must first exam the amplitude behavior under the limit of $l_T^2 \to 0$ for all the diffractive processes in the calculations using the two-gluon exchange model. From Eq. (13), we can see that the amplitude for the diffractive photoproduction of di-quark jet process is exact zero at $l_T^2 \to 0$. However, for the diffractive hadroproduction process $gp \to q\bar{q}p$ the amplitude for the first four diagrams is not exact zero in the limit $l_T^2 \to 0$ due to the inequality of the color factors between them. So, for $gp \to q\bar{q}p$ process there must be other diagrams in this order of

perturbative QCD calculation to cancel out the linear singularity which rises from the first four diagrams. The last five diagrams of Fig.2 are just for this purpose.

For example, the contributions from Diags.5 and 8 are

$$\mathrm{Im}\mathcal{A}^{58}(\pm,+,-)=\alpha_k^2(1+\alpha_k)\mathcal{N}\times\int\frac{d^2\vec{l}_T}{(l_T^2)^2}f(x',x'';l_T^2)\left(-\frac{1+\alpha_k}{4}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-(1+\alpha_k)\vec{l}_T)}{(\vec{k}_T-(1+\alpha_k)^2\vec{l}_T)^2}\right),\tag{15}$$

and the contributions from Diags.6 and 9 are

$$\mathrm{Im}\mathcal{A}^{69}(\pm,+,-)=\alpha_k^2(1+\alpha_k)\mathcal{N}\times\int\frac{d^2\vec{l}_T}{(l_T^2)^2}f(x',x'';l_T^2)\left(\frac{\alpha_k}{4}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-\alpha_k\vec{l}_T)}{(\vec{k}_T-\alpha_k^2\vec{l}_T)^2}\right),\tag{16}$$

and the contribution from Diag.7 is

$$\mathrm{Im}\mathcal{A}^{7}(\pm,+,-)=\alpha_k^2(1+\alpha_k)\mathcal{N}\times\int\frac{d^2\vec{l}_T}{(l_T^2)^2}f(x',x'';l_T^2)\left(\frac{1}{2}\frac{\vec{e}^{(\pm)}\cdot\vec{k}_T}{k_T^2}\right).\tag{17}$$

From the above results, we can see that the contributions from Diags.5-9 just cancel out the linear singularity which rises from the first four diagrams. Their total sum from the nine diagrams of Fig.2 is free of linear singularity now.

Finally, by adding up all of the nine diagrams of Fig.2, the imaginary parts of the amplitudes for the following helicity sets are,

$$\begin{aligned}\mathrm{Im}\mathcal{A}(\pm,+,-)&=\alpha_k^2(1+\alpha_k)\mathcal{N}\times\mathcal{T}^{(\pm)},\\ \mathrm{Im}\mathcal{A}(\pm,-,+)&=\alpha_k(1+\alpha_k)^2\mathcal{N}\times\mathcal{T}^{(\pm)},\end{aligned}\tag{18}$$

where

$$\begin{aligned}\mathcal{T}^{(\pm)}=\int\frac{d^2\vec{l}_T}{(l_T^2)^2}f(x',x'';l_T^2)\Big[&\left(\frac{1}{2}-\frac{2}{9}\right)\frac{\vec{e}^{(\pm)}\cdot\vec{k}_T}{k_T^2}-\frac{1}{36}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-\vec{l}_T)}{(\vec{k}_T-\vec{l}_T)^2}\\ &-\frac{1+\alpha_k}{4}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-(1+\alpha_k)\vec{l}_T)}{(\vec{k}_T-(1+\alpha_k)\vec{l}_T)^2}+\frac{\alpha_k}{4}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-\alpha_k\vec{l}_T)}{(\vec{k}_T-\alpha_k\vec{l}_T)^2}\Big].\end{aligned}\tag{19}$$

And the amplitudes for the other helicity sets are zero in the light quark jet production process. From the above results, we can see that in the integration of the amplitude the linear singularity from different diagrams are canceled out by each other, which will guarantee there is no linear singularity in the total sum.

Another feature of the above results for the amplitudes is the relation to the differential off-diagonal gluon distribution function $f(x',x'';l_T^2)$. To get the cross section for the diffractive process, we must perform the integration of Eq. (19). However, as mentioned above that there is no big difference between the off-diagonal gluon distribution function and the usual gluon distribution at small x, so we can simplify the integration of (19) by approximating the differential off-diagonal gluon distribution function $f(x',x'';l_T^2)$ by the usual diagonal differential gluon distribution function $f_g(x;l_T^2)$.

After integrating over the azimuth angle of $\vec{l}_T$, the integration $\mathcal{T}^{(\pm)}$ will then be

$$\mathcal{T}^{(\pm)} = \pi \frac{\vec{e}^{(\pm)} \cdot \vec{k}_T}{k_T^2} \mathcal{I}, \tag{20}$$

where

$$\mathcal{I} = \int \frac{dl_T^2}{(l_T^2)^2} f_g(x; l_T^2) [\frac{1}{36}(\frac{1}{2} - \frac{k_T^2 - l_T^2}{2|k_T^2 - l_T^2|}) + \frac{1+\alpha_k}{4}(\frac{1}{2} - \frac{k_T^2 - (1+\alpha_k) l_T^2}{2|k_T^2 - (1+\alpha_k) l_T^2|}) \\ - \frac{\alpha_k}{4}(\frac{1}{2} - \frac{k_T^2 - \alpha_k l_T^2}{2|k_T^2 - \alpha_k l_T^2|})]. \tag{21}$$

Comparing the above results with those of the photoproduction process $\gamma p \to q\bar{q}p$ [4,15], we find that the amplitude formula for the diffractive light quark jet hadroproduction process $gp \to q\bar{q}p$ is much more complicated. However, the basic structure of the amplitude, especially the expression for the integration $\mathcal{I}$ is similar to that for the photoproduction process. In the integration of (21), if $l_t^2 < k_T^2$ the first term of the integration over l_T^2 will be zero; if $l_t^2 < k_T^2/(1+\alpha_k)^2$ the second term will be zero; if $l_t^2 < k_T^2/\alpha_k^2$ the third term will be zero. So, the dominant regions contributing to the three integration terms are $l_t^2 \sim k_T^2$, $l_t^2 \sim k_T^2/(1+\alpha_k)^2$, and $l_t^2 \sim k_T^2/\alpha_k^2$ respectively. Approximately, by ignoring some evolution effects of the differential gluon distribution function $f_g(x; l_T^2)$ in the above dominant integration regions, we get the following results for the integration $\mathcal{I}$,

$$\mathcal{I} = \frac{1}{k_T^2}[\frac{1}{36} f_g(x; k_T^2) + \frac{(1+\alpha_k)^3}{4} f_g(x; \frac{k_T^2}{(1+\alpha_k)^2}) - \frac{\alpha_k^3}{4} f_g(x; \frac{k_T^2}{\alpha_k^2})]. \tag{22}$$

Obtained the formula for the integration $\mathcal{I}$, the amplitude squared for the partonic process $gp \to q\bar{q}p$ will be reduced to, after averaging over the spin and color degrees of freedom,

$$\overline{|\mathcal{A}|}^2 = \frac{9}{4} \alpha_s^3 (4\pi)^3 \pi^2 s^2 \frac{|\mathcal{I}|^2}{M_X^2}(1 - \frac{2k_T^2}{M_X^2}). \tag{23}$$

And the cross section for the partonic process $gp \to q\bar{q}p$ is

$$\frac{d\hat{\sigma}(gp \to q\bar{q}p)}{dt}|_{t=0} = \int_{M_X^4 > 4k_T^2} dM_X^2 dk_T^2 \frac{9\alpha_s^3 \pi^2}{8(M_X^2)^2} \frac{1}{\sqrt{1 - \frac{4k_T^2}{M_X^2}}} (1 - \frac{2k_T^2}{M_X^2}) |\mathcal{I}|^2. \tag{24}$$

The integral bound $M_X^2 > 4k_T^2$ above shows that the dominant contribution of the integration over M_X^2 comes from the region of $M_X^2 \sim 4k_T^2$. Using Eq. (4), this indicates that in this dominant region α_k is of order of 1. So, in the integration $\mathcal{I}$ the differential gluon distribution function $f_g(x; Q^2)$ of the three terms can approximately take their values at the same scale of $Q^2 = k_T^2$. That is, the integration $\mathcal{I}$ is then simplified to

$$\mathcal{I} \approx \frac{10M_X^2 - 27k_T^2}{36M_X^2} f_g(x; k_T^2). \tag{25}$$

Numerical calculations show that there is little difference (within 10% for $k_T > 5$ GeV) between the cross sections by using these two different parametrizations of $\mathcal{I}$, Eq. (22) and Eq. (25). So, in Sec.IV, we use Eqs. (24) and (25) to estimate the diffractive light quark jet production rate at the Fermilab Tevatron.

III. RECALCULATE THE HEAVY QUARK JET PRODUCTION USING THE HELICITY AMPLITUDE METHOD

For a crossing check, in this section we will recalculate the diffractive heavy quark jet production at hadron colliders by using the helicity amplitude method. In Ref. [13], we have calculated this process in the leading logarithmic approximation of QCD, where we expanded the amplitude in terms of l_T^2. Now, if we use the helicity amplitude method, we donot need to use the expansion method for the Γ_μ factor in Eq. (7) as in [13]. We can firstly calculate the amplitude explicitly by using the helicity amplitude method.

However, for the massive fermion, the amplitude formula is more complicated. Following Ref. [19], we first define the basic spinors $u_\pm(k_0)$ as,

$$
\begin{aligned}
u_+(k_0) &= \not{k}_1 u_-(k_0), \\
u_-(k_0)\bar{u}_-(k_0) &= \frac{1}{2}(1-\gamma_5)\not{k}_0,
\end{aligned}
\tag{26}
$$

where the momenta k_0 and k_1 satisfy the followoing relations,

$$
k_0 \cdot k_0 = 0, \quad k_1 \cdot k_1 = -1, \quad k_0 \cdot k_1 = 0. \tag{27}
$$

Using Eqs. (26) and (27), we can easily find that the spinor $u_+(k_0)$ satisfies

$$
u_+(k_0)\bar{u}_+(k_0) = \frac{1}{2}(1+\gamma_5)\not{k}_0. \tag{28}
$$

Provided the basic spinors, we then express any spinors $u(p_i)$ in terms of the basic ones, [19]

$$
u_\pm(p_i) = \frac{(\not{p}_i + m_i)u_\pm(k_0)}{\sqrt{2p_i \cdot k_0}}. \tag{29}
$$

It is easily checked that these spinors satisfy Dirac's equations. Now, for the massive fermions, the helicity amplitudes for the processes involving Dirac particles have the following general forms,

$$
\begin{aligned}
\bar{u}_+(p_f)Qv_-(p_i) &= \frac{Tr[Q(\not{p}_i - m_i)(1-\gamma_5)\not{k}_0(\not{p}_f + m_f)}{4\sqrt{(n\cdot p_i)(n\cdot p_f)}}, \\
\bar{u}_-(p_f)Qv_+(p_i) &= \frac{Tr[Q(\not{p}_i - m_i)(1+\gamma_5)\not{k}_0(\not{p}_f + m_f)}{4\sqrt{(n\cdot p_i)(n\cdot p_f)}}, \\
\bar{u}_+(p_f)Qv_+(p_i) &= \frac{Tr[Q(\not{p}_i - m_i)(1-\gamma_5)\not{k}_1\not{k}_0(\not{p}_f + m_f)}{4\sqrt{(n\cdot p_i)(n\cdot p_f)}}, \\
\bar{u}_-(p_f)Qv_-(p_i) &= \frac{Tr[Q(\not{p}_i - m_i)\not{k}_1(1-\gamma_5)\not{k}_0(\not{p}_f + m_f)}{4\sqrt{(n\cdot p_i)(n\cdot p_f)}},
\end{aligned}
\tag{30}
$$

where m_i and m_f are the masses for the momenta p_i and p_f respectively, where $p_i^2 = m_i^2$, $p_f^2 = m_f^2$. From the above equations, we can see that for the massless fermions ($m_i = m_f = 0$) the formula Eq. (30) will then turn back to the formula Eq. (12).

To calculate the imaginary part of the amplitude Eq. (7) for the partonic process $gp \to c\bar{c}p$, a convenient choice for the momenta k_0 and k_1 is,

$$k_0 = p, \qquad k_1 = e, \tag{31}$$

where the vector e is the polarization vector for the incident gluon defined in Eq. (11). Using the formula Eq. (30) and the above choice for the momenta k_0 and k_1, the helicity amplitudes for Eq. (7) will then be,

$$\begin{aligned}
\mathrm{Im}\mathcal{A}(\pm,+,-) &= \alpha_k^2(1+\alpha_k)\mathcal{N} \times \mathcal{T}_c^{(\pm)}, \\
\mathrm{Im}\mathcal{A}(\pm,-,+) &= \alpha_k(1+\alpha_k)^2\mathcal{N} \times \mathcal{T}_c^{(\pm)}, &(32)\\
\mathrm{Im}\mathcal{A}(\pm,+,+) = \mathrm{Im}\mathcal{A}(\pm,-,-) &= \alpha_k(1+\alpha_k)\mathcal{N} \times \frac{\pi m_c}{2}\mathcal{I}_c', &(33)
\end{aligned}$$

where $\mathcal{N}$ is the same as in Eq. (14), and the integrations $\mathcal{T}_c^{(\pm)}$ and $\mathcal{I}_c'$ are defined as

$$\begin{aligned}
\mathcal{T}_c^{(\pm)} = \int \frac{d^2\vec{l}_T}{(l_T^2)^2} f(x',x'';l_T^2)[(\frac{1}{2}-\frac{2}{9})\frac{\vec{e}^{(\pm)}\cdot\vec{k}_T}{k_T^2+m_c^2} - \frac{1}{36}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-\vec{l}_T)}{m_c^2+(\vec{k}_T-\vec{l}_T)^2} \\
-\frac{1+\alpha_k}{4}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-(1+\alpha_k)\vec{l}_T)}{m_c^2+(\vec{k}_T-(1+\alpha_k)\vec{l}_T)^2} + \frac{\alpha_k}{4}\frac{\vec{e}^{(\pm)}\cdot(\vec{k}_T-\alpha_k\vec{l}_T)}{m_c^2+(\vec{k}_T-\alpha_k\vec{l}_T)^2}], &\quad (34)\\
\mathcal{I}_c' = \frac{1}{\pi}\int \frac{d^2\vec{l}_T}{(l_T^2)^2} f(x',x'';l_T^2)[(\frac{1}{2}-\frac{2}{9})\frac{1}{k_T^2+m_c^2} - \frac{1}{36}\frac{1}{m_c^2+(\vec{k}_T-\vec{l}_T)^2} \\
-\frac{1+\alpha_k}{4}\frac{1}{m_c^2+(\vec{k}_T-(1+\alpha_k)\vec{l}_T)^2} + \frac{\alpha_k}{4}\frac{1}{m_c^2+(\vec{k}_T-\alpha_k\vec{l}_T)^2}]. &\quad (35)
\end{aligned}$$

If we approximate the differential off-diagonal gluon distribution function $f(x',x'';l_T^2)$ by the usual diagonal differential gluon distribution function $f_g(x;l_T^2)$, the above integrations will then be reduced to, after integrating over the azimuth angle of $\vec{l}_T$,

$$\mathcal{T}_c^{(\pm)} = \pi\vec{e}^{(\pm)}\cdot\vec{k}_T\mathcal{I}_c, \tag{36}$$

where

$$\begin{aligned}
\mathcal{I}_c = \int \frac{d^2\vec{l}_T}{(l_T^2)^2} f_g(x;l_T^2)[\frac{5}{18}\frac{1}{m_T^2} - \frac{5}{36}\frac{1}{k_T^2} - \frac{1}{72}\frac{k_T^2-m_c^2-l_T^2}{m_1^2} \\
-\frac{1+\alpha_k}{8}\frac{k_T^2-m_c^2-(1+\alpha_k)^2l_T^2}{m_2^2} + \frac{\alpha_k}{8}\frac{k_T^2-m_c^2-\alpha_k^2l_T^2}{m_3^2}], &\quad (37)
\end{aligned}$$

where

$$\begin{aligned}
m_T^2 = k_T^2+m_c^2, \quad m_1^2 = \sqrt{(m_t^2+l_T^2)^2-k_T^2l_T^2}, \quad m_2^2 = \sqrt{(m_t^2+(1+\alpha_k)^2l_T^2)^2-(1+\alpha_k)^2k_T^2l_T^2}, \\
m_3^2 = \sqrt{(m_t^2+\alpha_k^2l_T^2)^2-\alpha_k^2k_T^2l_T^2}, &\quad (38)
\end{aligned}$$

and

$$\mathcal{I}_c' = \int \frac{d^2\vec{l}_T}{(l_T^2)^2} f_g(x;l_T^2)[\frac{5}{18}\frac{1}{m_T^2} - \frac{1}{36}\frac{1}{m_1^2} - \frac{1+\alpha_k}{4}\frac{1}{m_2^2} + \frac{\alpha_k}{4}\frac{1}{m_3^2}]. \tag{39}$$

So, the amplitude squared for the partonic process $gp \to c\bar{c}p$ will then be reduced to, after averaging over the spin and color degrees of freedom,

$$\overline{|\mathcal{A}|}^2 = \frac{9}{4}\alpha_s^3(4\pi)^3\pi^2 s^2 \frac{m_T^2}{M_X^2}[(1-\frac{2k_T^2}{M_X^2})k_T^2|\mathcal{I}_c|^2 + m_c^2|\mathcal{I}_c'|^2]. \tag{40}$$

From the above results, we can see that the integrals of Eqs. (37) and (39) are proportional to $1/l_T^2$ in the limit of $l_T^2 \to 0$. That is to say that there exist large logarithmic contributions from the integration region of $1/R_N^2 \ll l_T^2 \ll m_T^2$ for the integration over l_T^2 as in Ref. [13]. So, we can expand the integrals of Eqs. (37) and (39) in terms of l_T^2 to get the leading logarithmic contribution to the amplitude. In the limit of $l_T^2 \to 0$, the parameters m_1^2, m_2^2 and m_3^2 scale as,

$$\begin{aligned}
\frac{1}{m_1^2} &\approx \frac{1}{m_T^2}[1-\frac{m_c^2-k_T^2}{m_T^2}\frac{l_T^2}{m_T^2}],\\
\frac{1}{m_2^2} &\approx \frac{1}{m_T^2}[1-\frac{m_c^2-k_T^2}{m_T^2}\frac{(1+\alpha_k)^2 l_T^2}{m_T^2}],\\
\frac{1}{m_3^2} &\approx \frac{1}{m_T^2}[1-\frac{m_c^2-k_T^2}{m_T^2}\frac{\alpha_k^2 l_T^2}{m_T^2}].
\end{aligned} \tag{41}$$

Under this approximation, the integrations Eqs. (37) and (39) will then be related to the integrated gluon distribution function $xg(x;Q^2)$,

$$\mathcal{I}_c \approx \frac{2m_c^2}{36(m_T^2)^3 M_X^2}(10M_X^2 - 27m_T^2)xg(x;m_T^2), \tag{42}$$

$$\mathcal{I}'_c \approx \frac{m_c^2-k_t^2}{36(m_T^2)^3 M_X^2}(10M_X^2 - 27m_T^2)xg(x;m_T^2). \tag{43}$$

Substituting the above results into Eq. (40), we can then reproduce the leading logarithmic approximation result for the diffractive charm jet production process at hadron colliders which has been calculated in [13].

IV. NUMERICAL RESULTS

Provided with the cross section formula (24) for the partonic process $gp \to q\bar{q}p$, we can calculate the cross section of the diffractive light quark jet production at the hadron level. However, as mentioned above, there exists nonfactorization effect caused by the spectator interactions in the hard diffractive processes in hadron collisions. Here, we use a suppression factor $\mathcal{F}_S$ to describe this nonfactorization effect in the hard diffractive processes at hadron colliders [8]. At the Tevatron, the value of $\mathcal{F}_S$ may be as small as $\mathcal{F}_S \approx 0.1$ [8,9]. That is to say, the total cross section of the diffractive processes at the Tevatron may be reduced down by an order of magnitude due to this nonfactorization effect. In the following numerical calculations, we adopt this suppression factor value to evaluate the diffractive production rate of light quark jet at the Fermilab Tevatron.

The numerical results of the diffractive light quark jet production at the Fermilab Tevatron are plotted in Fig.3 and Fig.4. In our calculations, the scales for the parton distribution

functions and the running coupling constant are both set to be $Q^2 = k_T^2$. For the parton distribution functions, we choose the GRV NLO set [20]. In Fig.3, we plot the differential cross section $d\sigma/dt|_{t=0}$ as a function of the lower bound of the transverse momentum of the light quark jet, $k_{T\min}$. This figure shows that the cross section is sensitive to the transverse momentum cut $k_{T\min}$. The differential cross section decreases over four orders of magnitude as $k_{T\min}$ increases from 5 GeV to 15 GeV.

It is interesting to compare the cross section of the diffractive light quark jet production with that of the diffractive charm quark jet production [13], which is also shown in Fig.3. The two curves in this figure show that the production rates of the light quark jet and the charm quark jet in the diffractive processes are in the same order of magnitude. However, we know that the cross section of diffractive heavy quark jet production is related to the integrated gluon distribution function, while the cross section of the light quark jet production is related to the differential gluon distribution function.

In Fig.4, we plot the differential cross section $d\sigma/dt|_{t=0}$ as a function of the lower bound of the momentum fraction of the proton carried by the incident gluon $x_{1\min}$, where we set $k_{T\min} = 5\ GeV$. This figure shows that the dominant contribution comes from the region of $x_1 \sim 10^{-2} - 10^{-1}$. This property is the same as that of the diffractive charm jet production at the Tevatron [13].

V. CONCLUSIONS

In this paper, we have calculated the diffractive light quark jet production at hadron colliders in perturbative QCD by using the two-gluon exchange model. We find that the production cross section is related to the squared of the differential gluon distribution function $\partial G(x; Q^2)/\partial lnQ^2$ at the scale of $Q^2 \sim k_T^2$, where k_T is the transverse momentum of the final state quark jet. For a crossing check, we also used the helicity amplitude method to calculate the diffractive charm jet production process at hadron colliders, by which we could reproduce the leading logarithmic approximation result for this process previously calculated in Ref. [13]. We have also compared the production rate of the light quark jet in the diffractive processes with that of the heavy quark jet production, and found that the production rates of these two processes are in the same order of magnitude.

As we know, the large transverse momentum dijet production in the diffractive processes at hadron colliders is important to study the diffractive mechanism and the nature of the Pomeron. The CDF collaboration at the Fermilab Tevatron have reported some results on this process [9]. In this paper, under the two-gluon exchange model in perturbative QCD we have calculated the light quark jet production in the diffractive processes. In a forthcoming paper, we will calculate the diffractive gluon jet production in hadron collisions by using the two-gluon exchange model, which will complete the calculations of the diffractive dijet production in perturbative QCD. These results then can be used to compare with the experimental measurements to test the valid of the perturbative QCD description of the diffractive processes at hadron colliders and may be used to study the diffractive mechanism and the factorization broken effects.

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Figure Captions

Fig.1. Sketch diagram for the diffractive quark jet production at hadron colliders in perturbative QCD.

Fig.2. The lowest order perturbative QCD diagrams for partonic process $gp \to q\bar{q}p$.

Fig.3. The differential cross section $d\sigma/dt|_{t=0}$ for the light quark jet production in the diffractive processes as a function of $k_{T\min}$ at the Fermilab Tevatron, where $k_{T\min}$ is the lower bound of the transverse momentum of the out going quark jet. For the charm quark jet production cross section formula, we take from Ref. [13] and set $m_c = 1.5\ GeV$.

Fig.4. The differential cross section $d\sigma/dt|_{t=0}$ for the light quark jet production as a function of $x_{1\min}$, where $x_{1\min}$ is the lower bound of x_1 in the integration of the cross section.

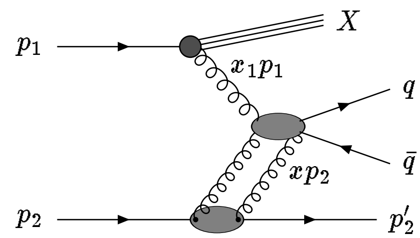


Fig.1

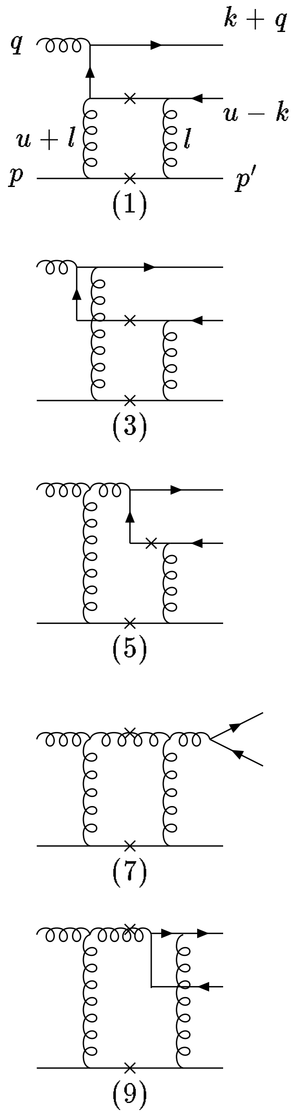
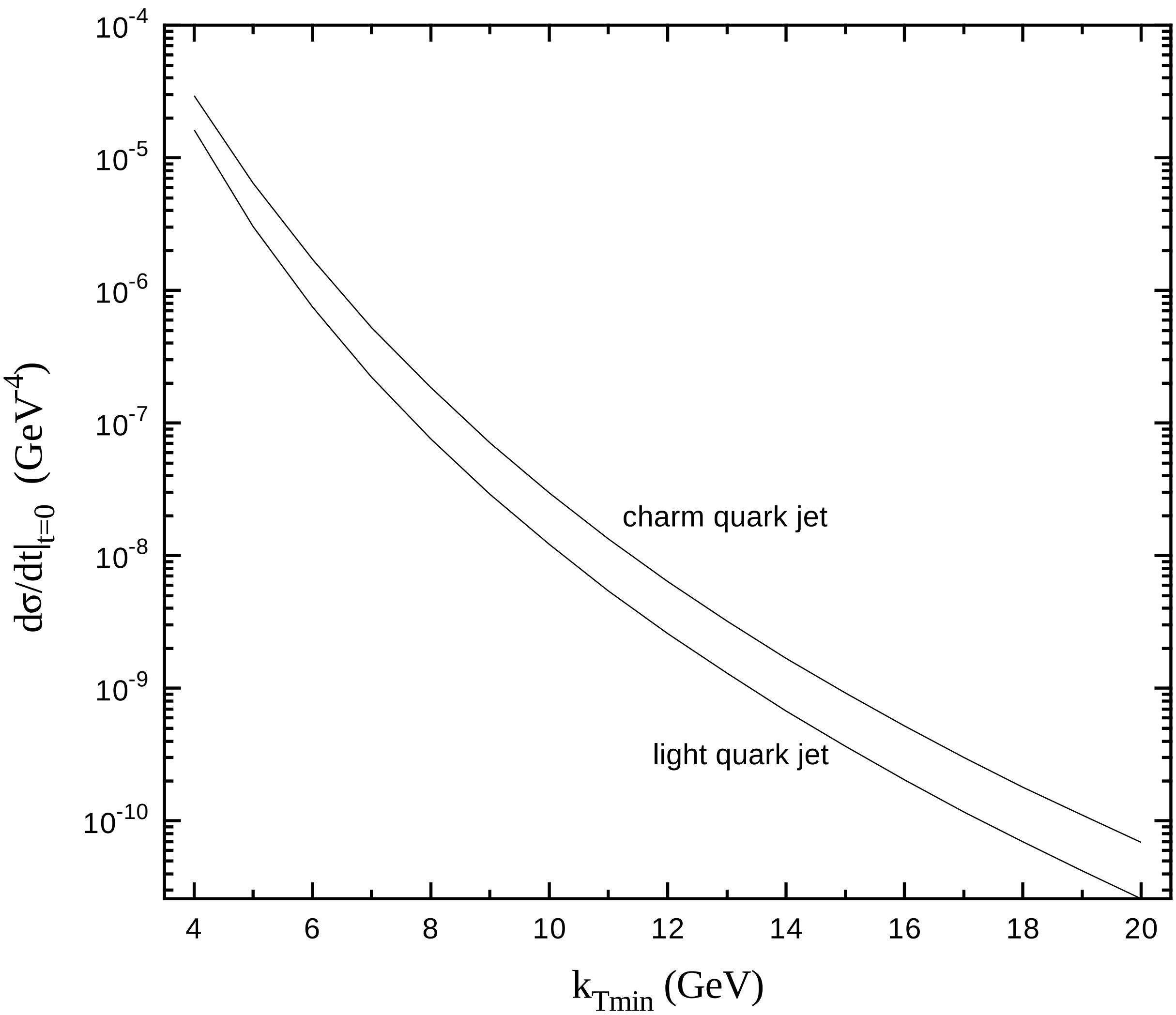


Fig.2



10^{-4}
10^{-5}
10^{-6}
10^{-7}
10^{-8}
10^{-9}
10^{-10}
$d\sigma/dt|_{t=0}$ (GeV^{-4})
charm quark jet
light quark jet
4
6
8
10
12
14
16
18
20
k_{Tmin} (GeV)

Fig.3

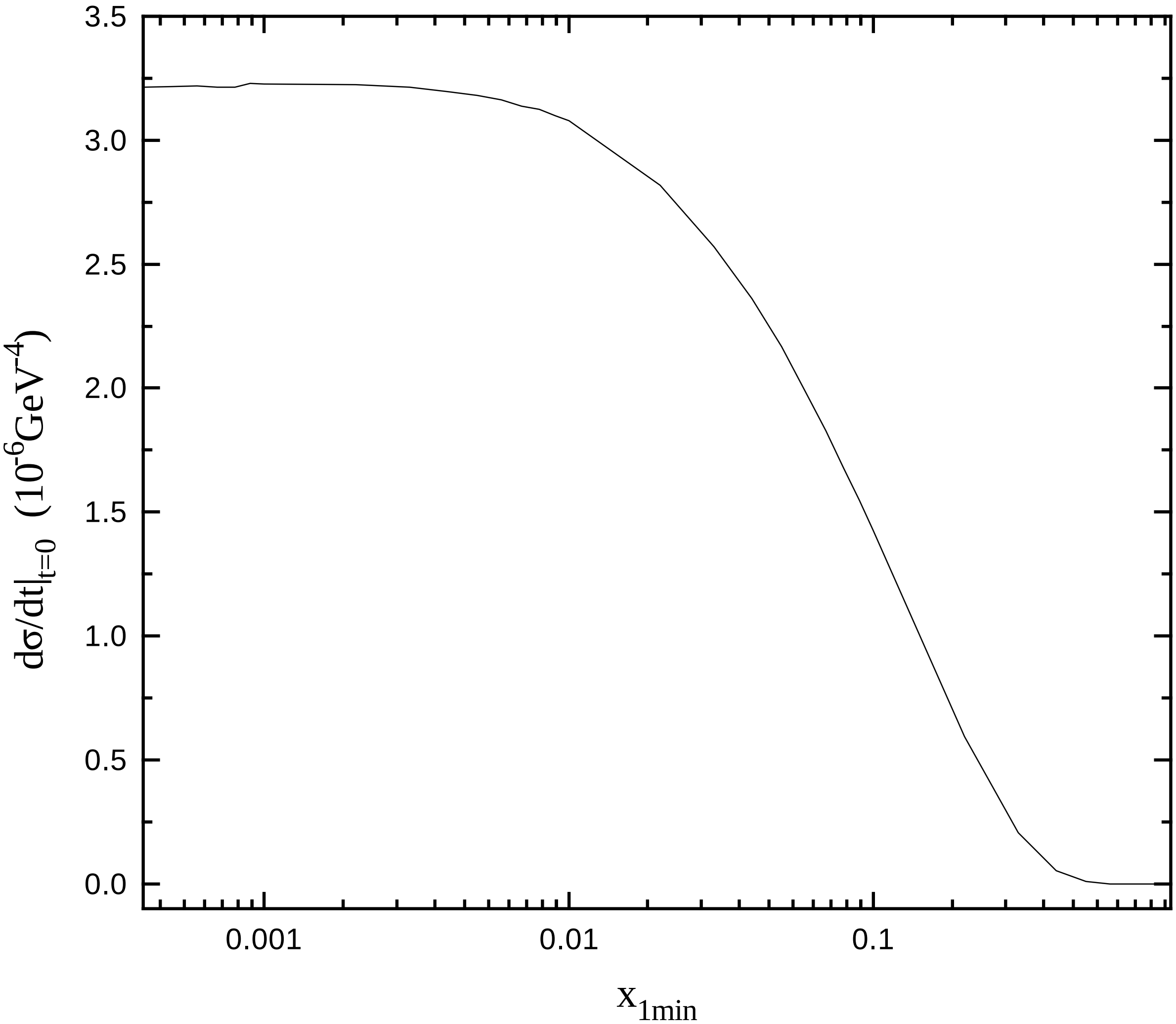


Fig.4